SUPERSYMMETRY AND THE SPONTANEOUS BREAKING OF A U(1) GAUGE FACTOR

E. Keith[1], Ernest Ma[2] and Biswarup Mukhopadhyaya[3]

Department of Physics

University of California

Riverside, CA 92521, USA

ABSTRACT

We investigate a supersymmetric theory with an extra U(1) gauge symmetry surviving down to low energies. The extra U(1) is assumed to originate from an E_6 grand unified theory (GUT). We show that if one assumes universal soft supersymmetry breaking scalar masses at the GUT scale, and requires the mass of the additional U(1) gauge boson to satisfy phenomenological bounds, then the conditions for electroweak symmetry breaking will provide stringent restrictions on the allowed parameter space of such a theory. We also determine the masses of standard as well as exotic sfermions and find that it is possible for the latter to be lighter than the former. An interesting specific observation is that a light stop is difficult to accommodate with an extra U(1) gauge symmetry.

[1] E-mail : evan@citrus.ucr.edu

[2] E-mail : ernestma@citrus.ucr.edu

[3] Permanent address: Mehta Research Institute, 10 Kasturba Gandhi Marg, Allahabad - 211 002, India.
E-mail: biswarup@mri.ernet.in

1 INTRODUCTION

Supersymmetry (SUSY) is perhaps the most widely pursued option in the quest for physics beyond the standard model [1]. Though the efforts in this direction are woven largely around the minimal supersymmetric standard model (MSSM), the importance of extending beyond the minimal case should also be recognized. There are at least two main reasons: (1) The mechanism for SUSY breaking is not yet understood, and a viable mechanism (although not necessarily the only one) consists in breaking SUSY softly at a very high scale, for example, through gravitational interactions [2]; (2) Recent measurements of the three coupling constants of the standard model have strongly suggested that they tend to unify at an energy scale of 10^{16} GeV in a supersymmetric scenario [3]. Thus there is enough motivation to think that if Nature is indeed supersymmetric, then SUSY is perhaps embedded in a framework of higher symmetry which, manifested at a high scale, ultimately leads to softly broken SUSY and all its desirable features at a scale of 1 TeV or below.

A class of candidate theories in this context consists of the E_6 grand unified theories (GUTs) which were originally motivated by superstrings. One generic feature of such theories is that they predict an extra U(1) symmetry surviving down to low energies [4]. Consequently, as we stand close to the exploration of phenomenology at the TeV scale, it is useful to know all possible effects of this extra U(1) . For example, the spontaneous breaking of this U(1) leads to an additional neutral gauge boson whose phenomenological implications are under close scrutiny today. It has also been suggested [5, 6] that such a scenario can lead to light singlet neutrinos which offer a possible reconciliation of the recent LSND data on neutrino oscillations [7] with the solar and atmospheric neutrino puzzles.

In this paper, we consider a general supersymmetric E_6 scenario, and discuss how the breakdown of the extra U(1) at the scale of about 1 TeV affects the sfermion masses in a SUSY model. With a universal scalar mass of size m_0 at the GUT scale, all the scalar masses at low energy are determined by the gaugino contributions to the running masses, as well as by the D-term contributions which arise from the reduction in rank of the gauge group at the SUSY breaking scale. In particular, the large vacuum expectation value (VEV) of an isosinglet scalar field, which in turn gives mass to the additional gauge boson, leads to

D-terms of substantial magnitude. We shall show that this affects the Higgs mass parameters in such a way that the model becomes quite restricted from the standpoint of electroweak symmetry breaking. In addition, it predicts certain qualitative features among the sfermion masses. The combinations of parameters which in general fulfill all constraints normally lead to sfermion masses beyond the reach of LEP-2, and perhaps beyond that of the Tevatron as well. Exotic sfermions lighter than ordinary ones can be envisioned in such a scenario. Also, the large D-terms make it difficult to accommodate a light stop, which is a rather interesting feature of such theories.

In Section 2 we discuss the salient features of a SUSY model with an extra U(1) from E_6. We also list there the general expressions for different scalar masses in this framework. In Section 3 the constraints that follow on model parameters are discussed in the light of these expressions. Section 4 contains some sample numerical results. There we also point out the improbability of a light stop in such a scenario. We conclude in Section 5. The appendix contains some general observations on the feasibility of breaking an extra U(1) symmetry while preserving SUSY.

2   GENERAL MODEL AND SCALAR MASSES

As has been previously stated, we consider a supersymmetric E_6 grand unified theory. The extra U(1)'s in E_6 can be described by the following decomposition:

$$E_6 \longrightarrow SO(10) \times U(1)_\psi \tag{1}$$

and

$$SO(10) \longrightarrow SU(5) \times U(1)_\chi . \tag{2}$$

A rank-5 low-energy theory can be obtained if a linear combination of $U(1)_\psi$ and $U(1)_\chi$ survives down to the TeV scale:

$$U(1)_\psi \times U(1)_\chi \longrightarrow U(1)_\alpha, \tag{3}$$

where the charge for the surviving extra U(1) is, in general,

$$Q_\alpha = Q_\psi \cos \alpha - Q_\chi \sin \alpha \tag{4}$$

The particle content we will choose to use here is three complete **27**'s of E_6 along with an extra pair of color-singlet $SU(2)_L$ doublet fields. The different quantum numbers, including the ψ- and χ-charges of the various superfields are found in Table 2 of Ref. [4]. However, we note that different conventions are followed by different authors. In one-loop order, our choice of field content leads to a unification of the gauge couplings at the same scale as in the MSSM. We also note that, due to the addition of exotic colored fields, α_s does not run in one-loop order. Of course, one could also have assumed that gauge-coupling unification comes from threshold corrections which could then occur at about the string scale.

The angle α is unspecified in general. The most common specific case is where E_6 breaking takes place via Wilson loops, leading to the η-model, where in our convention, $\alpha = \tan^{-1} \sqrt{3/5}$. Another example is a recently proposed model, henceforth to be referred to as the N-model, where light singlet neutrinos can be generated, which corresponds to $\alpha = -\tan^{-1} \sqrt{1/15}$. Details of this latter model can be found in Ref. [5].

The **27** of E_6 contains all the matter fields surviving down to the TeV scale. In addition to the particles belonging to the MSSM, each generation also has the isosinglet color-triplets h and h^c, and the isosinglet color-singlets ν^c and S, the latter being trivial under $U(1)_\chi$ as well. One can construct the model so as to break $U(1)_\alpha$ spontaneously with the VEV of the scalar component of S (as well as the two $SU(2)$ doublet Higgs VEVs):

$$\langle \tilde{S} \rangle = u \,, \ \langle H_1 \rangle = v_1 \,, \ \langle H_2 \rangle = v_2 \,, \tag{5}$$

where H_1 and H_2 are the two Higgs doublets giving masses to the down- and up-type quarks respectively, and $\tan \beta = v_2/v_1$. It is the singlet VEV u which sets the scale of $U(1)_\alpha$ breaking, to be constrained by experimental lower bounds on the mass of Z', the additional neutral gauge boson. There is in general mixing between Z and Z', and the physical states can be obtained by diagonalizing the mass matrix. In practice, assuming $g_\psi = g_\chi = g_Y$ from a GUT hypothesis, the Z' mass is approximately given by

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$$m_{Z'}^2 = \frac{4}{3}g_Y^2 u^2 \cos^2\alpha + O(v_1^2, v_2^2) \,. \tag{6}$$

Note that a value of α very close to $\pi/2$ will lead to an inadmissibly small Z' mass, and is hence disallowed for u in the TeV range.

Let us now consider the sfermion masses in this scenario. We assume that soft SUSY breaking scalar terms are generated via gravitational interactions with a hidden sector at the GUT scale M_G. We further make the conventional assumption that these soft breaking terms are flavor blind. Consequently, at the scale M_G, there is a universal mass m_0 for all scalar fields, a universal gaugino mass $M_i = m_{1/2}$, and all trilinear soft SUSY breaking scalar terms are parameterized by a universal massive parameter A_0. At the one-loop level, we have the gluino mass $m_{\tilde{g}} = m_{1/2}$ because the one-loop beta function vanishes for the strong coupling. The scalar masses at low energy are obtained by running the renormalization group equations (RGEs) down to the weak scale and including the D-term masses arising from both $U(1)_\alpha$ and electroweak symmetry breaking. We include F-term masses only when calculating the stop masses. The RGEs yield the following expressions for the scalar masses at the scale μ:

$$m_{H_2}^2 = m_0^2 + \sum_i C_i^{(H_2)} - \frac{3}{2}I \,, \tag{7}$$

$$m_{\tilde{Q}_3}^2 = m_0^2 + \sum_i C_i^{(Q_3)} - \frac{1}{2}I \,, \tag{8}$$

$$m_{\tilde{t}^c}^2 = m_0^2 + \sum_i C_i^{(t^c)} - I \,, \tag{9}$$

$$m_R^2 = m_0^2 + \sum_i C_i^{(R)} \,, \tag{10}$$

where R represents all other fields except those appearing in the trilinear superpotential term $\lambda_t H_2 \tilde{Q}_3 \tilde{t}^c$. We have assumed that only the top Yukawa coupling, which appears through

$$I = \frac{1}{4\pi^2} \int_t^{t_G} \lambda_t^2 \left(A_t^2 + m_{H_2}^2 + m_{\tilde{Q}_3}^2 + m_{\tilde{t}^c}^2 \right) dt \tag{11}$$

has a significant effect on the running of the scalar masses. In the above equation, the gaugino loop contribution for a scalar field F is given as [8]

$$C_i^{(F)} = c_2(F_i) \cdot \frac{1}{2\pi^2} \int_t^{t_G} g_i^2 M_i^2 dt \,, \tag{12}$$

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with $t = \ln(\mu/\text{GeV})$ and $t_G = \ln(\mu_G/\text{GeV})$. Both of these integrals can be solved analytically at the one-loop level [9].

At the weak scale and at the scale u, D-term contributions to these masses must be added to the previous expressions. Although in our numerical calculation we have included $O(v_1^2, v_2^2)$ contributions to these terms for completeness, here we list only the $O(u^2)$ terms, with a scalar field F receiving the D-term contribution of $u^2 g_\alpha^2 Q_\alpha(F) Q_\alpha(S)$ in the square of its mass:

$$D_{16(10)} = g_\alpha^2(\cos\alpha)u^2 \left(\frac{1}{6}\cos\alpha + \frac{1}{2\sqrt{15}}\sin\alpha \right), \tag{13}$$

$$D_{16(\bar{5})} = g_\alpha^2(\cos\alpha)u^2 \left(\frac{1}{6}\cos\alpha - \frac{3}{2\sqrt{15}}\sin\alpha \right), \tag{14}$$

$$D_{16(1)} = g_\alpha^2(\cos\alpha)u^2 \left(\frac{1}{6}\cos\alpha + \sqrt{\frac{5}{12}}\sin\alpha \right), \tag{15}$$

$$D_{10(5)} = g_\alpha^2(\cos\alpha)u^2 \left(-\frac{1}{3}\cos\alpha - \frac{1}{\sqrt{15}}\sin\alpha \right), \tag{16}$$

$$D_{10(\bar{5})} = g_\alpha^2(\cos\alpha)u^2 \left(-\frac{1}{3}\cos\alpha + \frac{1}{\sqrt{15}}\sin\alpha \right), \tag{17}$$

where the notation refers to the decomposition of E_6 under $SO(10)(SU(5))$ as $27 \rightarrow 16(10 + \bar{5} + 1) + 10(5 + \bar{5}) + 1(1)$ and the 16 contains all of the MSSM quark and lepton fields as well as ν^c. In our calculations, we assume $g_\alpha = g_Y$, where g_Y and g_α are both normalized in the grand unified group. This is a very good approximation because $b_\alpha \approx b_Y$ for the one-loop beta functions.

There are a few points to be noted here. First, the D-terms induced by the singlet VEV u can be much larger that the electroweak ones. This can cause large shifts in the scalar masses, in the positive or negative direction, depending on the sign of α. As we shall see in the next section, this implies nontrivial additional constraints on such a scenario. Secondly, the presence of three complete **27**'s slows down the evolution of the strong coupling α_s. In fact at the one-loop level, it does not evolve at all between the electroweak and GUT scales. As a result, the strong coupling stays stronger than in the MSSM case as the scale increases to the GUT scale. Thus the gluino loop contribution to the running masses turns out to be considerably higher than that in the MSSM. And finally, there is a new, but small, RGE contribution from the $U(1)_\alpha$ gaugino loop.

3 CONSTRAINTS FROM THE SCALAR MASS EXPRESSIONS

As can be clearly seen from the previous section, there are several parameters $(m_0, m_{\tilde{g}}, u,$ $\tan\beta, A_0, m_t, \alpha_s)$ determining the actual values of the scalar masses. However, certain qualitative features are seen among the mass terms derived with an additional broken U(1), which restrict the model considerably. We shall examine these constraints below.

The Z' mass is proportional to the singlet VEV u for any given value of α. The experimental lower limit on $m_{Z'}$ is model-dependent. For the η-model it is about 440 GeV [10], while a model-independent analysis of the Tevatron data, assuming the Z' to have the same coupling strength as the Z, sets a limit of nearly 700 GeV [11]. Setting $g_\psi = g_\chi = g_Y$ leads to a relatively weaker coupling for the Z', since it has no SU(2) coupling except through mixing. Thus in a general anaylsis it is safe to assume that the Z' mass must be at least 500 GeV. This immediately translates to a minimum value of u for any α.

The higher u is, the magnitudes of the D-terms are also correspondingly higher. The most important consequence of this will be the way the parameters in the Higgs sector are affected. A desirable (although technically not absolutely required) condition to be satisfied for electroweak symmetry breaking is

$$m_{H_2}^2 + \mu^2 < 0\,, \tag{18}$$

where the Higgsino mass parameter μ is given at tree level by

$$\mu^2 = -M_Z^2/2 + (m_{H_1}^2 - m_{H_2}^2 \tan^2\beta)/(\tan^2\beta - 1) \tag{19}$$

Also, there is the condition regarding the minimum value of the square of the pseudoscalar mass:

$$m_A^2 = m_{H_1}^2 + m_{H_2}^2 + 2\mu^2 > 0 \tag{20}$$

at tree level. In the above, we note that the experimental bound on m_A comes from the nonobservation of the decay $Z \to h + A$ which depends also on m_h. In the limit of large $\tan\beta$, $m_h = m_A$, hence $m_A > M_Z/2$ may be used. However, we will be conservative and use zero instead. Also, the existence of the extra U(1) at the TeV scale changes the two-Higgs-doublet

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structure at the electroweak scale. We have assumed these changes to be small enough so that the MSSM remains a valid approximation. We now rewrite Eqs. (18) and (20) as

$$-\frac{M_Z^2}{2} + \frac{\Delta m^2}{\tan^2\beta - 1} < 0 \tag{21}$$

and

$$-M_Z^2 - \frac{\Delta m^2}{\cos 2\beta} > 0\,, \tag{22}$$

where

$$\Delta m^2 = D_{10(\bar{5})} - D_{10(5)} + \frac{3}{2}I\,. \tag{23}$$

Figs. 1 and 2 show some sample parameter space of the N- and η-models, respectively, to indicate the types of constraints implied by these inequalities. Clearly, with a positive value of α it is rather difficult to satisfy the first condition (e.g. Fig. 2(a) for the η-model), while still having u large enough to be compatible with the experimental limits on $m_{Z'}$. The fact that I is always positive makes the constraint even stronger. The general expression for I is linear in both m_0^2 and $m_{\tilde{g}}^2$. In order that the damaging effect of I is less, one requires m_0 (as also the gaugino mass) to be on the lower side. Also, the large D-terms can be balanced only if $\tan\beta$ is sufficiently large. To satisfy the second inequality, however, a low value of m_0 or $m_{\tilde{g}}$ makes matters worse, as can be seen from Fig. 2(b). Note also that we need $\cos 2\beta$ here to be negative, hence $\tan\beta > 1$ is required. It turns out that in the η-model, the only way to satisfy the conditions is to have $\tan\beta \approx 8$ or higher. Models with negative values of α, like the N-model, are less tightly restricted because the D-term constrictions drive the first inequality in the right direction, favoring a large u. However, the pseudoscalar mass limit prevents u from being too large in this case also. On the whole, in such models it is possible to have a smaller value of $\tan\beta$ than, say, in the η-model, as the curves in Fig. 1(a) and 1(b) indicate.

A further constraint on the η-model comes from the nature of the two types of D-terms that contribute to the sfermion masses. Of them, $D_{16(10)}$ is always positive. $D_{16(\bar{5})}$, however, can be negative for a positive value of α. This in turn implies large negative shifts in $m_{\tilde{d}^c}^2$, $m_{\tilde{e}}^2$, and $m_{\tilde{\nu}}^2$. Of these, $m_{\tilde{d}^c}^2$ receives a substantial positive RGE contribution from the gluino loops. However, constraining the "left-handed" sneutrino masses to be above 45 GeV (the

limits from LEP-1) requires the universal scalar mass m_0 to be appropriately large, or requires the universal gaugino mass $m_{1/2} = m_{\tilde{g}}$ to be big enough to produce a compensating effect from the SU(2) and U(1) gaugino loops. This subjects the η-model to an added restriction, as is illustrated in Fig. 3. In the N-model, on the other hand, both $D_{16(\bar{5})}$ and $D_{16(10)}$ are positive, and thus this model is free from such a restriction. Another implication of this constraint is that the η-model, or any model with a positive α, is compatible with a no-scale supergravity scenario (where $m_0 = 0$) only if the universal gaugino mass is sufficiently high (at least about 300 GeV). We shall also see in the next section that the large positive values of $D_{16(10)}$, as well as the necessity of a large m_0, makes it difficult to have a light stop in this scenario.

4   SOME NUMERICAL RESULTS

In our calculations we have assumed the GUT scale to be 2×10^{16} GeV. The beta functions for the different couplings are determined using

$$b_1 = 3(3) + \frac{3}{5}, \quad b_2 = -6 + 3(3) + 1, \tag{24}$$

$$b_3 = -9 + 3(3), \quad b_N = 3(3) + \frac{2}{5}. \tag{25}$$

Also, we use $\alpha_s = 0.123$, $\alpha = 1/127.9$, and $\sin^2 \theta_W = 0.2317$ at the scale M_Z. We take the top mass to be 175 GeV.

The graphs presented in Fig. 4 show the masses of the sfermions belonging to the first two generations plotted against the Z' mass for some sample parameter space. The η- and N-models are chosen to demostrate general features in the case of positive and negative α respectively. The values for the other parameters chosen for the purpose correspond to regions in the parameter space which satisfy the constraints discussed in the previous section.

In both of the sample scenarios, the squark masses are predicted to lie approximately in the range 300-550 GeV, and the slepton masses in the range 100-450 GeV. This puts them almost beyond the search limits of the upgraded Tevatron, and above the limits of LEP-2. If they are discovered with lower masses, then that will render an additional U(1) factor somewhat unlikely, since that would make it extremely difficult to satisfy the constraints

mentioned above, for a generic α. On the other hand, the allowed range of the sfermion masses falls well within the discovery limits of the LHC. It should be mentioned here that if SUSY has to be broken at a scale not exceeding a TeV or so, then any additional Z' is quite likely to be discovered at the LHC together with the sfermions.

On the other hand, Fig. 5 shows that the exotic (SU(2) singlet) squarks can have relatively low masses. This is because h^c has both Q_ψ and Q_χ negative, while h has a positive Q_χ but an equally negative Q_ψ. As a result, the D-terms for them can reduce the scalar masses. Thus in such a scenario the exotic squarks, which should be produced on par with the ordinary ones at hadronic colliders, could be discovered *before* the ordinary ones. Of course, it should be remembered all along that there is a further uncertainty in the prediction of the exotic squark masses, since their F-term masses are not predicted by the model. In fact, if the Z' mass is large then it is imperative for the corresponding exotic quarks to be heavy so that one does not end up with an inadmissibly small value for $m_{\tilde{h}}^2$. Also, at least one of the three exotic quarks should be heavy for it to induce a sizable negative contribution by its Yukawa coupling to the square of the mass of $\tilde{S}$ to precipitate the spontaneous breaking of $U(1)_\alpha$ in the first place.

Finally, let us consider the stop masses in this framework. By virtue of the potentially large $\tilde{t}_L$-$\tilde{t}_R$ terms arising through trilinear couplings in the third generation, the stop mass matrix is given by

$$\mathcal{M}_{\text{stop}} = \begin{pmatrix} m_{\tilde{Q}_3}^2 + m_t^2 + D_{16(10)} & m_t(A_t + \mu \cot\beta) \\ m_t(A_t + \mu \cot\beta) & m_{\tilde{t}^c}^2 + m_t^2 + D_{16(10)} \end{pmatrix}, \tag{26}$$

where A_t is the trilinear scalar coupling for the third generation and we have neglected the $O(v_1^2, v_2^2)$ D-term mass contributions in the expression. Diagonalization of this mass matrix implies that one physical state $(m_{\tilde{t}_1})$ can be quite light (≤ 100 GeV) if the off-diagonal element happens to be large enough. This has attracted a lot of interest in recent times [12] because of two main reasons, namely (1) a partial solution to the R_b problem is offered by a light stop scenario [13], and (2) a light stop provides a better fit to $\alpha_s(m_Z^2)$ [14]. Both of these require the stop to be lighter than, or at most, about 100 GeV. It has been found that in the MSSM framework, a stop in this mass range can indeed exist without contradicting the currently available data on top decay from the Tevatron, although with accumulating

luminosity one can restrict it further. While LEP-1 imposes a lower limit of 45 GeV on the stop mass, the only other constraint on it is from a search for the direct production of stop pairs, which closes a window in the range 65-88 GeV. Thus a stop lighter than about 100 GeV is an object of active interest.

In our scenario, one unavoidable feature is the presence of the new D-term in the diagonal mass terms $m^2_{\tilde{t}_{L,R}}$. This is the term $D_{16(10)}$ which is always positive. The effect of such a term is to boost both the eigenvalues of the mass matrix. The only way in which a light stop could still be envisioned is in cases where the off-diagonal terms are also very large. This would require the magnitude of either A_t or μ to be large compared to the diagonal mass parameters. The former, however, is constrained by conditions arising from color and charge invariance of the vacuum, as well as by flavor-changing neutral current suppression [15]. Our numerical estimates show that the lowest (≈ 100 GeV) possible value of $m_{\tilde{t}_1}$ necessitates $A_0 \sim 1$ TeV where the above conditions cannot be satisfied. On the other hand, a large magnitude of μ requires the universal scalar mass m_0 or $m_{\tilde{g}}$ to be correspondingly large, a situation that would simultaneously increase the diagonal terms as well. Thus the general conclusion is that a light stop of the type envisioned in the MSSM is rather unlikely when there is an additional U(1) symmetry broken at such a scale that the additional gauge boson has a phenomenologically permissible mass.

In Fig. 6 we show our expectations of the lighter stop mass in regions of the parameter space where m_0 and $m_{\tilde{g}}$ are close to their minimal possible values. Although the N-model allows smaller values for the lightest stop than the η-model does, due to possible smaller values of m_0 and $m_{\tilde{g}}$, even here we do not find a stop mass lighter than 120 GeV when we scan the viable parameter space. The different predicted values of μ and $f = |\mu/u|$ (the coupling of the superpotential term SH_1H_2 in the model) for the same parameter space as in Fig. 6 are plotted in Fig. 7.

5   SUMMARY AND CONCLUSIONS

We have studied the scalar mass patterns predicted in a SUSY scenario embedded in a general E_6 grand unified theory, with an extra U(1) gauge symmetry surviving at low energy.

Our observation is that, if one assumes scalar mass universality at the GUT scale, then the conditions for electroweak symmetry breaking as well as a phenomenologically viable pseudoscalar mass puts a general model under constraints. Such constraints are especially tight for the η-model, as also for models where the angle α is negative. As far as the standard sfermion masses are concerned, they are expected to be in a somewhat higher range than in the MSSM case, although they can still be within the discovery limits of the LHC. This scenario also allows the interesting possibility of lighter exotic squarks. And lastly, it is difficult to accommodate a light stop in such a picture.

6   ACKNOWLEDGEMENTS

APPENDIX

In this paper we have discussed the case of a spontaneously broken U(1) gauge symmetry together with the explicit soft breaking of supersymmetry. However, it is also important to consider the possibility of preserving SUSY while breaking the U(1) gauge symmetry, or more generally, while reducing the rank of the gauge symmetry, *e.g.* from SO(10) to SU(5). Suppose we want to break G to G' such that the rank of G' is one less than the rank of G. Then there must be an U(1) gauge factor such that

$$G \supset G' \times U(1), \tag{27}$$

and the scalar superfield which does it must transform under this $U(1)$. If there is only one such superfield, it is clearly impossible to have a superpotential invariant under G and hence no spontaneous breaking of G is possible without also breaking supersymmetry.

Consider now two scalar superfields, ϕ_1 and ϕ_2, transforming oppositely under G and hence also under $U(1)$, then the superpotential is

$$W = \mu \phi_1 \phi_2. \tag{28}$$

The supersymmetric scalar potential is then

$$V = |\mu \phi_1|^2 + |\mu \phi_2|^2 + V_D, \tag{29}$$

where $V_D \geq 0$ comes from the gauge sector. The supersymmetric minimum of this V clearly does not break the gauge symmetry.

We now add a singlet superfield χ, *i.e.* one which is trivial under G. The superpotential is then

$$W = \mu \phi_1 \phi_2 + f \phi_1 \phi_2 \chi + r\chi + \frac{1}{2}M\chi^2 + \frac{1}{3}h\chi^3, \tag{30}$$

and the supersymmetric scalar potential becomes

$$V = |\mu \phi_1 + f \phi_1 \chi|^2 + |\mu \phi_2 + f \phi_2 \chi|^2 + |f \phi_1 \phi_2 + r + M\chi + h\chi^2|^2 + V_D. \tag{31}$$

A supersymmetric minimum of V is obtained for

$$\langle \chi \rangle = -\frac{\mu}{f}, \quad \langle \phi_1 \phi_2 \rangle = -\frac{r}{f} + \frac{M\mu}{f^2} - \frac{h\mu^2}{f^3}, \tag{32}$$

and $|\langle\phi_1\rangle| = |\langle\phi_2\rangle|$ from $V_D = 0$.

Consider the possible effect on scalar masses due to the so-called D-terms. The latter are proportional to $|\langle\phi_1\rangle|^2 - |\langle\phi_2\rangle|^2$ and have thus no effect as long as supersymmetry is maintained. In the presence of soft SUSY breaking by universal scalar masses at a chosen scale, the scalar-mass parameters at the symmetry breaking scale would differ by how they couple to the other particles of the theory in the evolution of these couplings away from the chosen scale. In the minimal supersymmetric standard model (MSSM), the two Higgs scalar doublets evolve differently because one couples to the t quark (with its corresponding large coupling) and the other does not.

Recently there have been discussions in the literature[16] regarding the contribution of D-terms to the masses of the MSSM sfermions from the reduction in rank of the gauge group at a very high scale. This is of course possible, but in order to calculate the deviations for a particular pattern of symmetry breaking, an implicit assumption of universal scalar masses at some higher scale (for example the Planck scale m_{pl}) must be made. Possible (but unknown) differences in the interactions of ϕ_1 and ϕ_2 with other fields between m_{pl} and $\langle\phi_{1,2}\rangle$ would then generate these extra D-terms whose magnitudes are however not guaranteed to be observable. In this paper, our D-terms are directly related to the mass of the Z' just as the usual electroweak D-terms are related to M_Z, and are thus subject to direct experimental verification.

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Figure captions

Fig. 1 : (a): $(m_{H_2}^2 + \mu^2)/\text{GeV}^2$ vs. $M_{Z'}/\text{GeV}$ for $\alpha = N$ with $m_{\tilde{g}} = 200$ GeV, $A_0 = 0$, and $\tan\beta = 4$.

(b): m_A^2/GeV^2 vs. $M_{Z'}/\text{GeV}$ for $\alpha = N$ with $m_{\tilde{g}} = 200$ GeV, $A_0 = 0$, and $\tan\beta = 4$.

In both case, the curves are labeled by the number m_0/GeV.

Fig. 2 : (a): $(m_{H_2}^2 + \mu^2)/\text{GeV}^2$ vs. $M_{Z'}/\text{GeV}$ for $\alpha = \eta$ with $A_0 = 0$, and $\tan\beta = 8$.

(b): m_A^2/GeV^2 vs. $M_{Z'}/\text{GeV}$ for $\alpha = \eta$ with $A_0 = 0$, and $\tan\beta = 8$.

In both cases, the curves are labeled by the pair of numbers $(m_0/\text{GeV}, m_{\tilde{g}}/\text{GeV})$.

Fig. 3 : Lower bound on m_0/GeV vs. $M_{Z'}/\text{GeV}$ for $\alpha = \eta$ from $m_{\tilde{\nu}_{e_L}} > 45$ GeV. The different curves represent gluino masses from 150 GeV to 400 GeV at 50 GeV intervals.

Fig. 4 : (a): $m_{\tilde{f}}/\text{GeV}$ vs. $M_{Z'}/\text{GeV}$ for $\alpha = N$ with $m_{\tilde{g}} = 200$ GeV, $m_0 = 0$ GeV, $A_0 = 0$, and $\tan\beta = 4$. In descending order on the left-hand end of the graph, the curves represent the masses of $\tilde{d}_R$, $\tilde{d}_L$, $\tilde{u}_L$, $\tilde{u}_R$, $\tilde{e}_L$, $\tilde{\nu}_{e_L}$, and $\tilde{e}_R$.

(b): $m_{\tilde{f}}/\text{GeV}$ vs. $M_{Z'}/\text{GeV}$ for $\alpha = \eta$ with $m_{\tilde{g}} = 200$ GeV, $m_0 = 200$ GeV, $A_0 = 0$, and $\tan\beta = 8$. In descending order on the left-hand end of the graph, the curves represent the masses of $\tilde{d}_L$, $\tilde{u}_L$, $\tilde{u}_R$, $\tilde{d}_R$, $\tilde{e}_R$, $\tilde{e}_L$, and $\tilde{\nu}_{e_L}$.

Fig. 5 : (a): (exotic sfermion masses)/GeV vs. $M_{Z'}/\text{GeV}$ for $\alpha = N$ with $m_{\tilde{g}} = 200$ GeV, $m_0 = 0$ GeV, $A_0 = 0$, $\tan\beta = 4$, and $m_h = m_{h^c} = 400$ GeV.

(b): (exotic sfermion masses)/GeV vs. $M_{Z'}/\text{GeV}$ for $\alpha = \eta$ with $m_{\tilde{g}} = 180$ GeV, $m_0 = 200$ GeV, $A_0 = 0$, $\tan\beta = 8$, and $m_h = m_{h^c} = 200$ GeV.

Fig. 6 : (a): $m_{\tilde{t}_1}/\text{GeV}$ vs. $M_{Z'}/\text{GeV}$ for $\alpha = N$ with $m_{\tilde{g}} = 200$ GeV, $m_0 = 0$ GeV, $A_0 = 0$, and $\tan\beta = 4$.

(b): $m_{\tilde{t}_1}/\text{GeV}$ vs. $M_{Z'}/\text{GeV}$ for $\alpha = \eta$ with $m_{\tilde{g}} = 200$ GeV, $m_0 = 200$ GeV, $A_0 = 0$, and $\tan\beta = 8$.

Fig. 7 : (a): $|\mu|/\text{GeV}$ vs. $M_{Z'}/\text{GeV}$. The $\alpha = N$ case has $m_{\tilde{g}} = 200$ GeV, $m_0 = 0$ GeV, $A_0 = 0$, and $\tan\beta = 4$. The $\alpha = \eta$ case has $m_{\tilde{g}} = 200$ GeV, $m_0 = 200$ GeV, $A_0 = 0$, and $\tan\beta = 8$.

(b): $f = |\mu/u|$ vs. $M_{Z'}/\text{GeV}$. The $\alpha = N$ case has $m_{\tilde{g}} = 200$ GeV, $m_0 = 0$ GeV, $A_0 = 0$, and $\tan\beta = 4$. The $\alpha = \eta$ case has $m_{\tilde{g}} = 200$ GeV, $m_0 = 200$ GeV, $A_0 = 0$, and $\tan\beta = 8$.

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